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A scheme for ultrahigh-speed all-optical switching based on semiconductor heterostructures was explored theoretically. The approach was anticipated to be applicable to systems operating in the 10 to 100 gigabit/sec regime. Our role was to simulate the carrier dynamics in semiconductor heterostructures under ultrafast (picosecond and subpicosecond) laser excitation. Pulse characteristics were forund to coherently populate and depopulate the semiconductor heterostructure of carriers on the picosecond timescale.

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High-Speed All Optical Switching

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Abstract

A scheme for ultrahigh-speed all-optical switching based on semiconductor heterostructures was explored theoretically. The approach was anticipated to be applicable to systems operating in the 10 to 100 gigabit/sec regime. Our role was to simulate the carrier dynamics in semiconductor heterostructures under ultrafast (picosecond and subpicosecond) laser excitation. Pulse characteristics were forund to coherently populate and depopulate the semiconductor heterostructure of carriers on the picosecond timescale.

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The transmission of data through optical-fiber communications systems is ultimately limited by the optical bandwidth of silica fiber well in excess of 1 terahertz. Current systems, however, are limited to a small faction of that bandwidth due to limitations associated with modulating and demodulating optical data streams. One approach to high bandwidth optical communications systems is optical time-division multiplexing (OTDM) several data channels. Each individual channel is modulated at rates consistent with high-frequency electronics. Putting the channels together (multiplexing) and then teasing them apart (demultiplexing) requires optical switches that are as fast as the aggregate bit rate. Although nonlinear fiber loop devices have been demonstrated with the requisite speed, a purley semiconductor-based approach is desired to facilitate integration of optical switching devices with other components.

We explored the use of specially shaped optical pulses to devise a possible switching scheme that could operate at the requisite rates in a purely semiconductor-based structure. Namely, by coherently driving carrier populations up and then back down, we planned to demonstrate theoretically a route to optical switching that would meet the requisite 100 gigabit/sec rates, but for pipeline architectures. The quintessential phenomenon, well known from atomic physics, is Rabi flopping. The main question is: to what extent can Rabi flopping associated with the creation and coherent destruction of excitons be controlled? Whereas in atoms the approach is based on the optical Bloch equations (which describe the coherent coupling of a two-level system to an electromagnetic field), for semiconductors the appropriate description is based on the semiconductor Bloch equations which describe the evolution of the interband polarization and the electron and hole populations coupled to the electromagnetic field.

Our most important result is that one can use specially temporally shaped picosecond optical pulses to coherently control carrier populations in quantum wells. In particular, strongly chirped and frequency-modulated pulses were found to lead to pronounced population trapping effects and excitonic Rabi oscillations. Despite the fact that these effects were strong, it is not clear that they are very useful directly for optical switching, since following trapping, the carrier population goes into free-carrier states (i.e., not back down to the crystal ground state).

Due to the termination of this program well before the goals could be reasonably met, portions of this work were carried out under other funding subsequent to termination. Nevertheless, it should be noted that the funding from AFOSR was essential to begin work in this area in three key ways. (1) It funded Dr. Harshawardhan Wanare as a post-doctoral research fellow for less than a year. Dr. Wanare is an expert at coherent atom/field interactions; we required his expertise to carry over the relevant physics to the semiconductor domain. (2) Due to the widespread and well recognized difficulties finding post-doctoral research fellows in optics, it was not until April 1998 that Dr. Wanare could be hired. Nevertheless, it became evident that the numerical problems we were about to tackle were considerably more difficult than we had anticipated. We therefore used funds from this grant to purchase computer equipment (see attached memos). This equipment was required to perform computations supported by this grant [D. S. Citrin and W. Harshawardhan, "Terahertz sideband generation in quantum wells viewed as resonant photon tunneling through a time-dependent barrier: An exactly solvable model," Phys. Rev. B **60**, 1759 (1999); W. Harshawardhan, S. Hughes, and D. S. Citrin, "Ultrafast coherence effects in semiconductor quantum wells using modulated optical fields," Centenniel Meeting of the American Physical Society, Talk IC15.07] as well as to carry out work initially anticipated to be funded, but ultimately not funded by this grant but supported by other sources as an emergency measure [S. Hughes, W. Harshawardhan, and D. S. Citrin, "Excitonic state trapping and quasiadiabatic population transfer in a two-band semiconductor," Phys. Rev. B **60**, 15523 (1999)]. (3) The grant partially supported two Ph.D. students (Michael Feise and Alex Maslov) who contributed to the numerical computations.

In addition, this work has had a spin-off which we have pursued subsequently under different funding. In S. Hughes and D. S. Citrin, "Extremely tunable terahertz emission: Coherent population flopping in a dc-biased quantum well," *Optics in 1999*, Optics and Photonics News **10** (12), 44 (1999). In this work we showed theoretically that it may be possible to exploit excitonic Rabi flopping in semiconductor quantum wells as a source of coherent transient electromagnetic radiation in the 1-10 THz regime. This is exciting since there are few convenient sources of coherent transient radiation in this portion of the spectrum.

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AFOSR, 2447-0020
Grant F496209710535

May 12, 1999

To Whom It May Concern:

This letter is to justify the expenditure of 45 % of grant F496209710535 for equipment when no equipment was included in the original budget.

The start date of this three-year project was 11/1/97. The budget included the salary for a postdoctoral research associate for the full three years. However, it was not until 4/27/98 that I was able to hire a post-doctoral research associate to work on the project. When it had become evident that there would be a delay in hiring the post-doctoral research associate and that the computational complexity of the project exceeded our expectations, I decided to use some of the funds that were originally budgeted for salaries (plus associated fringe benefits and indirect costs) to purchase computer equipment. I believed that this was essential to meet the long-term goals of the project. This included three Digital workstations (DEC DPW-500) as well as associated peripherals. The original requisition for the equipment dates from 11 December 1997 whereas the purchase order is dated 26 February 1998. The equipment itself was received during June of 1998, but was not paid for until much later because the vendor did not fulfill all the terms of the agreement until then. These machines were used by me, my graduate students, and the postdoctoral research associate (after 4/27/98) to carry out nonlinear optical-pulse propagation in low-dimensional semiconductor structures (microcavities), as had been outlined in my proposal. A technical report was submitted detailing the progress made during the first year.

However, on August 20, 1999, I received notification from Dr. Alan Craig that the project would not be continued past the first year, which ended 10/31/98 (see copy of the notification attached). This happened because the Innovative Science and Technology directorate at BMDO was "disestablished" and was not due to any failure on my part. In spite of the fact that it is currently unfunded, I have continued to work on the project using the computer equipment which was purchased by the project. A list of publications submitted and conference presentations given that were supported by this project is available.

Please do not hesitate to contact me for further information.

David S. Citrin
Assistant Professor of Physics and Materials Science

From: alan.craig@afosr.af.mil
Date: Thu Aug 20 15:20:30 1998
To: <citrin@wsu.edu>
Subject: BMDO funds eliminated

Dear David,

The Innovative Science and Technology directorate at BMDO has been disestablished. (Lou Lome, the IS&T champion of efforts to explore high bandwidth optical networks for distributed computation with whom I have worked for many years, has been reassigned to a Simulation and Modeling group.) As a result, funds previously allocated to agents acting for BMDO to pursue technologies that support its goals have disappeared and programs dependent on them have been discontinued. Unfortunately, the termination of the IS&T function affects the continuation of your grant through my office. I have struggled to ameliorate the impact of this disaster, but am unable to discover any means to circumvent the budgetary difficulties. As a result, the remaining option increments on your grant will not be exercised.

Your investigation into the theory of coherent electronic states in semiconductors has only just begun in this program, so I apologize for the brief support. Together we gambled that the BMDO IS&T program would survive; now we endure the awkwardness of losing the bet. Only because I hope so fervently that perceptions in this arena will yield new technology did the bet make sense; now, in a program that supports primarily experimental work, it can no longer be justified. With time and persistence this area should garner renewed support. Perhaps your efforts over the past year have sufficiently established this field that it will attract recognition and investment at agencies with more expansive vision.

Thank you for your participation in the AFOSR/BMDO program.

Sincerely,
Alan

August 13, 1998

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WASHINGTON STATE UNIVERSITY	BUDGET STATEMENT BUDGET TITLE	BUDGET	DATES INCLUDED FROM	TO	PAGE
	PHYSICS CMR SPONSORED PROJECTS	2447	11-01-97	03-31-1999	1

PROGRAM 11F PROJECT 0020 High-Speed All-Optical Switching — 100% OF PERIOD ELAPSED
PROJECT PERIOD: 11-01-97 TO 10-31-98
PRIN. INVESTIGATOR: CITRIN, D

(1) PROJECT SUMMARY TO DATE BY OBJECT	PROJECT BUDGET	EXPENDED TO DATE	OUTSTANDING ENCUMBRANCES	BALANCE	PCT. USED
00 SALARIES	19,026.78	19,026.78		.00	100
03 GOODS AND SERVICES	374.56	374.56		.00	100
07 EMPLOYEE BENEFITS	4,463.62	4,463.62		.00	100
13 FACILITIES & ADMINISTRATION COST	10,116.04	10,116.04		.00	100
16 NON-CAPITALIZED EQUIPMENT	30,335.00	30,335.00		.00	100
PROJECT TOTAL.	64,316.00	64,316.00		.00	100

(2) PROJECT SUMMARY TO DATE	ESTIMATE	REVENUE	VARIANCE	PCT.
99 REVENUE	64,316.00	64,316.00	.00	100

(3) PROJECT SUMMARY TO DATE BY SUBOBJECT	PROJECT BUDGET	EXPENDED TO DATE	OUTSTANDING ENCUMBRANCES	BALANCE	PCT. USED
00-AF FACULTY		14,424.28			
00-AH GRADUATE ASSISTANTS		4,602.50			
TOTAL SALARIES.	19,026.78	19,026.78		.00	100
03-AC INSTRUCTION/LAB/MEDICAL SUPPLIES		59.51			
03-BW INTERDEPARTMENT SUPPLIES & SERVICES		279.00			
03-DS ROADRUNNER TOLLS		45.83			
03-HC REPRODUCTION		5.86			
03-WV GRANT OVERDRAFT REIMBURSEMENT		15.64-			
TOTAL GOODS AND SERVICES.	374.56	374.56		.00	100
07-DB TIAA/CREF 5%		699.96			
07-DG MEDICAL AID & WORKER'S COMPENSATION		158.85			
07-MA UNEMPLOYMENT INSURANCE		72.08			
07-MB STATE HEALTH INSURANCE		2,001.68			
07-MC GRADUATE STUDENT HEALTH INSURANCE		146.19			
07-QT QUALIFIED TUITION REDUCTION		1,384.86			
TOTAL EMPLOYEE BENEFITS	4,463.62	4,463.62		.00	100
13-DA DOMESTIC (ON-CAMPUS)		10,116.04			
TOTAL FACILITIES & ADMINISTRATION	10,116.04	10,116.04		.00	100
16-EB COMPUTING EQUIPMENT		30,335.00			
TOTAL NON-CAPITALIZED EQUIPMENT .	30,335.00	30,335.00		.00	100

Terahertz sideband generation in quantum wells viewed as resonant photon tunneling through a time-dependent barrier: An exactly solvable model

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(Received 7 April 1999)

The propagation of light through quantum wells in the vicinity of an excitonic resonance subjected to a THz electric field is considered. The treatment makes explicit the relationship between THz sideband generation and resonant photon tunneling through a time-dependent resonant-tunneling structure. A method for the coherent control of excitons in quantum wells is also suggested. [S0163-1829(99)11927-1]

The basis for much work on transport through harmonically varying time-dependent potentials dates back to the work of Tien and Gordon[1] on tunneling between superconducting films. Subsequent developments in semiconductor-heterostructure growth and split-gate devices largely account for the renaissance of interest in this area. A parallel development is the free-electron laser. The modulation via microwave and THz fields of transport properties of quantum-point contacts,[2] single-electron devices,[3] and resonant-tunneling structures[4,5] has been studied. Effort has also been put into observing THz sidebands (TS's) in the optical spectra of semiconductor quantum wells (QW's).[6] Despite substantial theoretical progress in the area of tunneling through time-dependent potentials, experimental results have been scant. And apart from the common role of the low-frequency field (microwave, THz) in inducing mutiphoton transitions in the system,[7] an *explicit* elucidation of the connection between resonant tunneling through time-dependent potentials and TS generation (TSG) was hitherto unavailable.[8] The connection though may prove valuable since resonant-tunneling devices based on semiconductor heterostructures introduce numerous complications prohibiting the controlled variation of a single parameter. These difficulties have hampered a detailed analysis of transport experiments through time-dependent potentials. Typically, experimental results have been modeled using Tien-Gordon theory which only captures the most basic effects.[7,9] Recently, however, exciting work has been published showing that the quantum mechanical phase of electrons can be measured in *tour de force* transport experiments[10,11] lending urgency to the exploitation of optical analogues where phase measurements are routine. Optical analogues of time-dependent transport problems therefore hold out the intriguing prospect of realizing a number of effects that have been theoretically predicted. Already, considerable experimental work on superluminality and its impact on issues concerning tunneling times has been carried out.[12]

In this study we present a theoretical treatment of light propagation through a QW in which the exciton energy, linewidth, and oscillator strength are modulated with bandwidths in the 100-GHz to 5-THz range (henceforth denoted THz field), such as provided by excitonic Stark effects induced by a THz field $\mathbf{F}(t)$ polarized in the QW plane $\hat{\mathbf{x}}$ or the quantum-confined Stark effect with $\mathbf{F}(t)$ parallel to the growth direction $\hat{\mathbf{z}}$. The formalism employed makes *explicit* the close relationship between TSG and resonant tunneling of photons through the QW in the presence of a time-dependent dielectric constant. As is well known,[13] the wave and Schrödinger equations can be mapped from one to the other via $(\omega/c)n(\omega,\mathbf{r})\leftrightarrow\{(2m/\hbar^2)[E(\omega)-V(\mathbf{r})]\}^{1/2}$, respectively, where the symbols here have their usual meanings. Previous studies of TSG (Ref. 6) have avoided the details of the propagation effects which bring to the fore the close analogy between the two classes of phenomena.

We consider linear optical propagation although the THz field may be strong. Thus, we are concerned with the low-density regime in which excitation-induced dephasing and saturation may be neglected. Let an optical pulse [amplitude $I(t)$] be incident normally on the QW in the vicinity of an excitonic resonance. The reflected $R(t)$ and transmitted $T(t)$ optical fields are given by[14]

$$\begin{bmatrix} T(t) \\ 0 \end{bmatrix} = \begin{bmatrix} \int_{-\infty}^{t} dt'\,[\delta(t-t')+\chi(t,t')] & \int_{-\infty}^{t} dt'\,\chi(t,t') \\ -\int_{-\infty}^{t} dt'\,\chi(t,t') & \int_{-\infty}^{t} dt'\,[\delta(t-t')-\chi(t,t')] \end{bmatrix} \begin{bmatrix} I(t') \\ R(t') \end{bmatrix}, \tag{1}$$

with $\chi(t,t')$ the optical susceptibility (to be discussed below) and the integrals operate on everything to the right.[14] We eliminate $R(t)$ from Eq. (1) to obtain

$$T(t)=I(t)+\int_{-\infty}^{t}dt'\,\chi(t,t')T(t'). \tag{2}$$

If $\dot{\mathbf{F}}=\mathbf{0}$, then $\chi(t,t')=\chi(t-t')$, and Eq. (2) is easily solved in the frequency domain; however, for $\dot{\mathbf{F}}\neq\mathbf{0}$, the susceptibility no longer depends simply on $t-t'$, but explicitly on the two time variables. Nevertheless, we shall see that even in this case, under certain conditions an explicit solution obtains.

The optical susceptibility $\chi(t,t')$ in the presence of $\mathbf{F}(t)$ can be treated as in Ref. 15 or solved for numerically;[16] however, as we show, a simple model can be considered which leads to a closed-form solution of Eq. (2). To motivate a physically reasonable *ansatz* for $\chi(t,t')$, first consider the case $\dot{\mathbf{F}}=\mathbf{0}$. For a single resonance (e.g., retaining the $1s$ exciton, but neglecting the other bound states and continuum) we have[14] $\chi(t-t')=\Gamma e^{-i(\varepsilon-i\gamma)(t-t')}$ where Γ is the exciton radiative width,[17] ε is the frequency of the exciton measured from the crystal ground state, and γ is the nonradiative linewidth. The parameters ε, γ, and Γ are for $\dot{\mathbf{F}}(t)=\mathbf{0}$. All other interband resonances of the QW are neglected and the line shape is assumed for convenience to be Lorentzian. If the Hamiltonian governing the interband electronic excitations of the QW depends on time through a THz-frequency electric field, but with frequency content sufficiently below any intraband resonance, we can account for the effects of the THz field through an adiabatic Stark shift, line broadening, and modification of the oscillator strength. (For the parameters investigated here, ultrahigh-quality ZnSe or GaN QW's might be the best candidates due to their high exciton binding energies.) We can thus use the ansatz $\chi(t,t')=G(t)G^*(t')$ with $G(t)=\Gamma^{1/2}(t)e^{-i\int^t dt'[\varepsilon(t')-i\gamma(t')]}$ where $\varepsilon(t)$, $\Gamma(t)$, and $\gamma(t)$ all depend on time due to the influence of $\mathbf{F}$.[18] The transfer-matrix method is a formulation of the scattering approach to such problems and is well known to follow from the tunneling Hamiltonian for the appropriate structure.[18] The reflectivity and transmission coeffeicients are *equivalent* to the Green function of the appropriate tunneling Hamiltonian.

Equation (2) can be converted into a differential equation. The solution is

$$T(t)=I(t)+\int_{-\infty}^{t}dt'\,\chi(t,t')I(t')e^{-\int_{t'}^{t}dt''\Gamma(t'')}. \tag{3}$$

This defines the response function $\delta(t-t')+\tilde{\chi}(t,t')$ with $\tilde{\chi}(t,t')=g(t)g^*(t')$, and

$$g(t)=\Gamma^{1/2}(t)e^{-i\int^t dt'\{\varepsilon(t')-i[\gamma(t')+\Gamma(t')]\}}.$$

Using $T=I+R$ [cf. Eq. (1)], we have $R(t)=\int_{-\infty}^{t}dt'\,\tilde{\chi}(t,t')I(t')$.

In the following we assume the THz field $\mathbf{F}(t)=\mathbf{F}_{\rm dc}+\mathbf{F}_{\rm ac}\cos\Omega t$ is applied to the QW with $\mathbf{F}_{\rm dc}$ (which may be zero) a dc bias and $\mathbf{F}_{\rm dc}\|\mathbf{F}_{\rm ac}$. We write the modulated parameters as $\varepsilon(t)=\varepsilon_0+\varepsilon_1\cos\zeta t$, $\gamma(t)=\gamma_0+\gamma_1\cos\zeta t$, and $\Gamma(t)=\Gamma_0-\Gamma_1\cos\zeta t$ where $\zeta=\Omega$ if $F_{\rm dc}\gg F_{\rm ac}$ or $\zeta=2\Omega$ if $F_{\rm dc}=0$. Typical values may be found in Refs. 19. We set $\tilde{\varepsilon}_i=\varepsilon_i-i\gamma_i$ and $\tilde{\varepsilon}_i'=\tilde{\varepsilon}_i-i\Gamma_i$ with $i=0,1$. If $F_{\rm dc}=0$, Stark effects are quadratic and the parameters are modulated at frequency 2Ω. This leads to even TS's of Ω as is expected for a system with inversion symmetry. If $F_{\rm dc}\neq 0$, the system lacks inversion symmetry and both even and odd TS's occur.

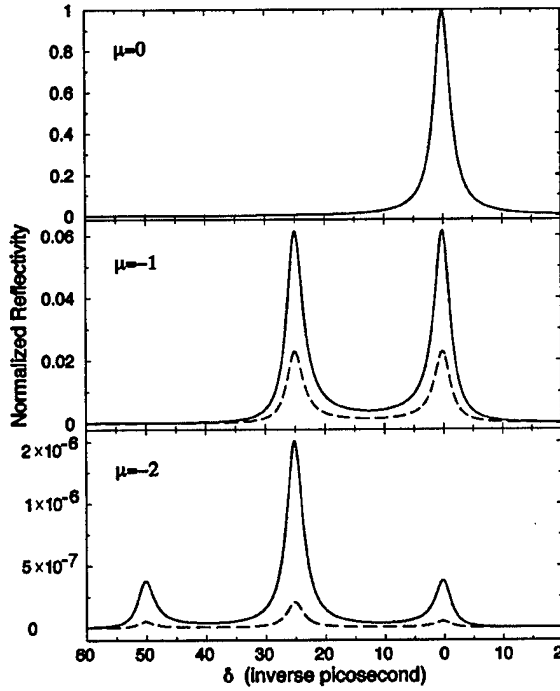

FIG. 1. Reflectivity $|R_\omega(\omega+\mu\zeta)|^2$ into harmonic μ normalized to peak reflectivity of the unmodulated fundamental $(\Gamma/\Gamma')^2=1.7\%$ as a function of detuning δ from exciton line center. The parameters are $\varepsilon_1=2.5$ (solid) and 1.52 $\rm ps^{-1}$ (dashed); in all cases $\gamma_0=1.52\ \rm ps^{-1}$, $\Gamma=0.1\ \rm ps^{-1}$, and $\gamma_1=0$.

Henceforth, we assume $\Gamma_1=0$. Let $I(t)=e^{-i\omega t}$. Then

$$T(t)=T_\omega(t)=e^{-i\omega t}+\Gamma e^{-i\omega t}\int_{-\infty}^{t}dt'\,e^{-i(\tilde{\varepsilon}_0'-\omega)(t-t')}e^{-i(\tilde{\varepsilon}_1/\zeta)[\sin\zeta t-\sin\zeta t']}. \tag{4}$$

Set $\tau=t-t'$, $\bar{t}=(t+t'/2)$, and $\sin\zeta t-\sin\zeta t'=2\sin(\zeta\tau/2)\cos\zeta\bar{t}$. Fourier transform $T_\omega(\omega')=\int_0^{2\pi/\zeta}dt\,e^{i\omega' t}T_\omega(t)$ to obtain $T_\omega(\omega')=(2\pi/\zeta)[(\omega-\tilde{\varepsilon}_0)/(\omega-\tilde{\varepsilon}_0')\delta_{\omega,\omega'}+\mathcal{K}_\mu(\omega)\delta_{\omega-\omega',\mu\zeta}]$ with

$$\mathcal{K}_\mu(\omega)=2i\Gamma\Delta\sum_{k=1}^{\infty}{}'\frac{1}{\Delta^2-(k\zeta/2)^2}\times J_{(k+\mu)/2}\left(\frac{\tilde{\varepsilon}_1 k}{2\Delta}\right)J_{(k-\mu)/2}\left(\frac{\tilde{\varepsilon}_1 k}{2\Delta}\right), \tag{5}$$

$\bar{\omega}=(\omega+\omega')/2$ the average of the incoming and outgoing frequencies, μ an integer, $\Delta=\Delta_\mu(\omega)=\bar{\omega}-\tilde{\varepsilon}_0'=\omega-\mu\zeta/2-\tilde{\varepsilon}_0'$ the complex detuning between the average frequency and the radiatively renormalized excitonic resonance, and

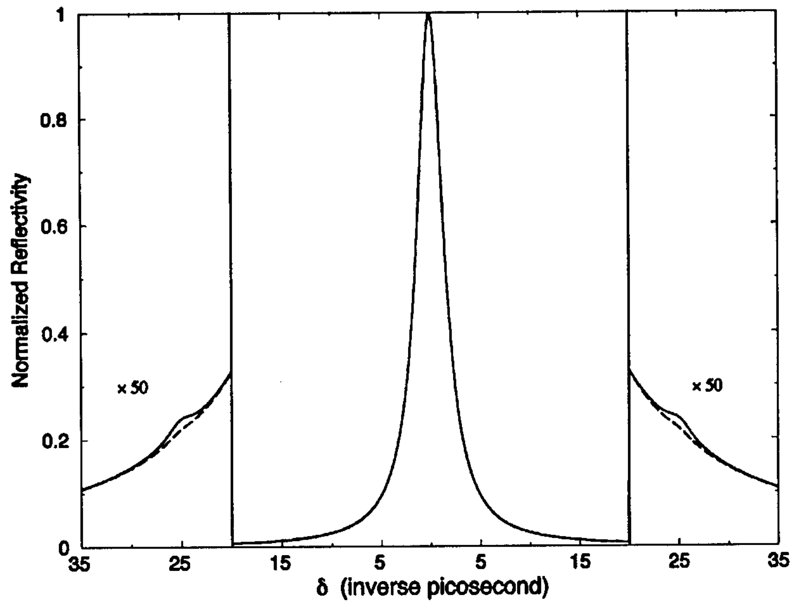

FIG. 2. Total reflected probability $R_{\rm tot}(\varepsilon)$ for the same parameters as Fig. 1 (solid). Note that for clarity, the region $|\delta|>20$ ps^{-1} is multiplied by a factor of 50.

$\Sigma_{k=1}^{\infty\prime}$ the sum restricted to integers k of the same parity as μ. $R_\omega(\omega')$ is obtained from the previous formula by replacing $(\omega-\tilde{\varepsilon}_0)/(\omega-\tilde{\varepsilon}_0')\delta_{\omega,\omega'}$ by $i\Gamma/(\omega-\tilde{\varepsilon}_0')\delta_{\omega,\omega'}$. The first term leads to the unmodulated result,[14] $\varepsilon_1,\gamma_1=0$. The other term gives the TS's as well as the modified transmission of the fundamental. Finally, for $\tilde{\varepsilon}_1=\varepsilon_1$ ($\gamma_1=0$), direct integration gives for the total reflected probability per unit time $R_{\rm tot}(\omega)=(\zeta/2\pi)\int_0^{2\pi/\zeta}dt\,|R_\omega(t)|^2=(\zeta/2\pi)\times(\Gamma/\Gamma')\mathrm{Re}\,R_\omega(\omega)$ (optical theorem) where $\Gamma'=\Gamma+\gamma_0$.

The foregoing expressions for R and T are compact; however, it is useful to reexpress the Kapteyn series for $\mathcal{K}_\mu(\omega)$ in Eq. (5) as a power series in $z=\tilde{\varepsilon}_1/\zeta$ (Ref. 20) to identify the various multi-THz-photon processes that contribute to a given TS:

$$\mathcal{K}_\mu(\omega)+i\frac{\Gamma}{\Delta}\delta_{\mu,0}=\begin{cases} i\Gamma\Delta\displaystyle\sum_{n=|\mu|/2}^{\infty}\left(\frac{\tilde{\varepsilon}_1}{2}\right)^{2n}\binom{2n}{|\mu|}\prod_{j=0}^{n}[\Delta^2-(j\zeta)^2]^{-1}, & \mu \text{ even}\\ i\Gamma\displaystyle\sum_{n=(|\mu|+1)/2}^{\infty}\left(\frac{\tilde{\varepsilon}_1}{2}\right)^{2n-1}\binom{2n-1}{|\mu|}\prod_{j=1}^{n}\left[\Delta^2-\left(\frac{2j-1}{2}\zeta\right)^2\right]^{-1}, & \mu \text{ odd.}\end{cases} \quad (6)$$

Series (5) and (6) converge in $|ze^{\sqrt{1-z^2}}/(1+\sqrt{1-z^2})|<1$ and for arbitrary $2\Delta/\zeta$ not an even (odd) integer.[20] The terms to a given order in $\tilde{\varepsilon}_1$ summed over μ are the nonlinear susceptibilities. For fixed μ, the term of degree p in $\tilde{\varepsilon}_1$ describes pth-order multi-THz-photon processes contributing to the TS. This in turn allows one to write the quantum mechanical transmission amplitude as a sum of nonlinear susceptibilities, as is commonly carried out in optics. The expansion (6) in terms of susceptibilities applies *mutatis mutandis* to the transport case, and to our knowledge was hitherto unknown.

Depending on whether $\mathbf{F}$ is polarized in the $\hat{\mathbf{x}}$ or $\hat{\mathbf{z}}$ direction, the dominant effect will differ; for $\mathbf{F}\|\hat{\mathbf{x}}$, the time dependence of $\gamma(t)$ will dominate through the ionization rate, while if $\mathbf{F}\|\hat{\mathbf{z}}$, $\varepsilon(t)$ will play a larger role.[19] Of course, $\varepsilon(t)$, $\gamma(t)$, and $\Gamma(t)$ are closely interlinked (see comments in Ref. 5), but for our purposes we can treat them as parameters determined by a microscopic theory external to our model.

To illustrate the different types of behaviors expected in these cases, we present results for the normalized reflectivity $(\Gamma'/\Gamma)^2|R_\omega(\omega+\mu\zeta)|^2$ [$(\Gamma/\Gamma')^2=0.017$ is the peak unmodulated reflectivity] of a high-quality QW in Fig. 1. We take $\gamma_1=0$ and consider the two cases 2.5 (solid) and 1.52 ps^{-1} (dashed) for $\mu=0$, 1, and 2. The results are plotted as a function of $\delta=\mathrm{Re}\,\Delta$. [In all cases, $\Gamma=0.1$ ps^{-1}, $\gamma_0=1.52$ ps^{-1} ($\hbar\gamma_0=1$ meV), and $\zeta/2\pi=4$ THz.] This gives the reflected probability into the μth TS as a function of incident frequency. The normalized reflectivity for $\mu=1$ is in the few % range while that of $\mu=2$ in the 10^{-4}% range for the parameters. Figure 2 shows $R_{\rm tot}(\omega)$ for the same parameters. The relative maximum strength of the first TS with

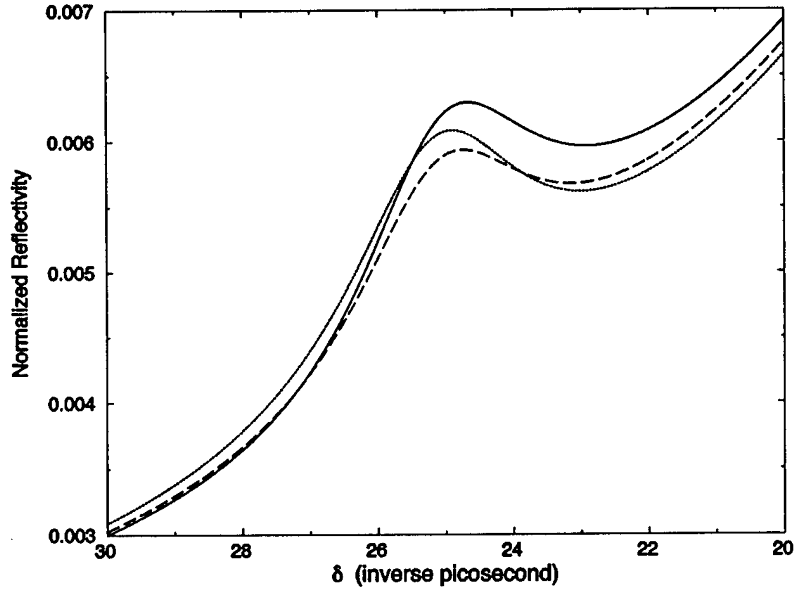

FIG. 3. Total reflected probability for $\tilde{\varepsilon}_1=\sqrt{2}\pi$ ps^{-1} (dotted), $(1+i)\pi$ ps^{-1} (dashed), and $i\sqrt{2}\pi$ ps^{-1} (solid). The parameters are otherwise those of Fig. 1.

respect to the unmodulated response is $\sim 14.5\%$ for $\varepsilon_1=2.5$ ps^{-1}. Figure 3 shows $R_{\rm tot}(\omega)$ for $\tilde{\varepsilon}_1=\sqrt{2}\pi$ ps^{-1} (dotted), $(1+i)\pi$ ps^{-1} (dashed), and $i\sqrt{2}\pi$ ps^{-1} (solid) with the other parameters the same as in Fig. 1 to show the effect of modulating the line center versus the linewidth. In all cases $|\tilde{\varepsilon}_1|=\sqrt{2}\pi$ ps^{-1}.

We turn our attention to time-domain effects. Specifically, we describe a new technique for coherent control of excitons in QW's using phase-locked pulse pairs. Previous schemes[21,22] utilize the relative phase between two identical but time-delayed optical pulses coincident with a narrow optical resonance in a QW. If the two pulses are in phase, the interband polarizations excited by each are in phase, and as a result ideally four times the population of excitons associated with a single pulse are generated; if, however, the pulses are π out of phase, the interference is destructive and the QW is coherently depopulated of excitons following the passage of the second pulse. Or at least that is the idea; dephasing, nonlinearities, and inhomogeneous broadening all conspire to make this effect less marked than one might hope.[22] In the scheme proposed here, two identical phase-locked time-delayed pulses are also used, but rather than modulate the relative phase of the two pulses, we introduce a phase shift in the polarization induced by the first optical pulse by a half-cycle THz pulse incident on the QW between the two optical pulses. Set $I(t)=I_0(t+t_0/2)+I_0(t-t_0/2)$ where $I_0(t)=A(t)e^{-i\varepsilon_0 t}$; $A(t)$ is the slowly varying envelope of a single constituent pulse of duration $\tau_0 \ll t_0 \ll \gamma_0^{-1}$, the time delay between pulses. Also let $\varepsilon(t)=\varepsilon_0+\varepsilon_1 B(t)$ and $\gamma(t)=\gamma_0$. $B(t)$ is a peaked function of time centered at $t=0$ whose duration is much less than t_0. For example, $B(t)$ might be due to a photoconductively generated half-cycle THz pulse propagating down a transmission structure deposited on the QW. For $t \gg t_0$,

$$T(t)=\Gamma\int_{-\infty}^{\infty}dt'\,e^{-i\tilde{\varepsilon}_0'(t-t')}e^{-i\varepsilon_1\int_{t'}^{t}dt''B(t'')}I(t')$$

$$=\Gamma e^{-i\tilde{\varepsilon}_0' t}\left[e^{-i\varepsilon_1\int_{-\infty}^{\infty}dt''B(t'')}+1\right]\int_{-\infty}^{\infty}dt'\,A(t') \quad (7)$$

and $R(t)=T(t)$. If the THz pulse is chosen so that $\varepsilon_1\int_{-\infty}^{\infty}dt''\,B(t'')=(2n+1)\pi$, destructive interference between the polarizations excited by the two optical pulses occurs and thus coherent depopulation of the excitons ensues; if the action integral is $2n\pi$, then constructive interference ensues.

As an example, if the half-cycle THz pulse is of 0.5 ps duration, we need $\hbar\varepsilon_1 \approx 2$ meV to achieve a π phase shift. This is feasible using available THz techniques, though carrying out this experiment may be difficult. One needs very high-quality QW's at low temperatures to ensure narrow excitonic lines. In addition, the THz transmission structure should be designed with $\mathbf{F}(t)\|\hat{\mathbf{z}}$ to suppress modulation of the dephasing; however, this can be compensated for by decreasing the amplitude of the second optical pulse in the sequence. Finally, the desirable effect will be degraded if there is temporal overlap of $B(t)$ and $A(t\pm t_0/2)$, thus requiring clean optical and THz wave forms. We note that the present scheme in no way circumvents the central difficulties with coherent depopulation of the exciton level, namely, inhomogeneous broadening and excitation-induced dephasing.

To conclude, we present an exactly solvable model for TSG by and photon propagation through THz-modulated QW's. Because of the mapping between the electromagnetic wave equation and the Schrödinger equation, light propagation through THz-modulated QW's will be a fruitful way to explore resonant tunneling through time-dependent potentials. We believe the area of optical analogues of time-dependent quantum mechanical transport problems—as pioneered by the work reviewed in Ref. 12—is in its infancy, and in particular, exploitation of optical pulse propagation through QW's shows great promise to advance the field mterially. In particular, we have explored the effects of modulating the exciton linewidth and energy on the TS's appearing in the optical spectra. We have also proposed an alternative scheme for carrying out two-pulse coherent control of excitons in QW's.

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[1] P. K. Tien and J. P. Gordon, Phys. Rev. **129**, 647 (1963).

[2] S. Feng and Q. Hu, Phys. Rev. B **48**, 5354 (1993); Q. Hu, Appl. Phys. Lett. **62**, 837 (1993); R. A. Wyss *et al.*, *ibid.* **63**, 1522 (1993).

[3] J. D. White and M. Wagner, Phys. Rev. B **48**, 2799 (1993).

[4] A.-P. Jauho *et al.*, Phys. Rev. B **50**, 5528 (1994); M. Wagner, Phys. Status Solidi B **204**, 382 (1997); W. Cai *et al.*, Phys. Rev. Lett. **65**, 104 (1990).

[5] N. S. Wingreen *et al.*, Phys. Rev. B **48**, 8487 (1993); N. S. Wingreen *et al.*, Phys. Rev. Lett. **61**, 1396 (1988).

[6] K. B. Nordstrom *et al.*, Phys. Rev. Lett. **81**, 457 (1998); K. Johnsen and A.-P. Jauho, Phys. Status Solidi B **204**, 55 (1997); W. Pötz, Phys. Rev. Lett. **79**, 3262 (1997).

[7] B. J. Keay *et al.*, Phys. Rev. Lett. **75**, 4098 (1995); B. J. Keay *et al.*, *ibid.* **75**, 4102 (1995).

[8] Harmonic generation in tunnel currents is discussed in A. I. Larkin and K. A. Matveev, Zh. Éksp. Teor. Fiz. **93**, 1030 (1987) [Sov. Phys. JETP **66**, 580 (1987)].

[9] P. S. S. Guimaraes *et al.*, Phys. Rev. Lett. **70**, 3792 (1993).

[10] R. Schuster *et al.*, Nature (London) **385**, 417 (1997).

[11] A. Yacoby *et al.*, Phys. Rev. Lett. **74**, 4047 (1995).

[12] R. Y. Chiao and A. M. Steinberg, in *Progress in Optics*, edited by E. Wolf (North-Holland, Amsterdam, 1997), Vol. 37, p. 345.

[13] D. A. Park, *Classical Dynamics and Its Quantum Analogues,* Lecture Notes in Physics Vol. 110 (Springer-Verlag, Berlin, 1979).

[14] L. C. Andreani, Phys. Lett. A **192**, 99 (1994); D. S. Citrin, Phys. Rev. B **50**, 5497 (1994).

[15] D. S. Citrin, Appl. Phys. Lett. **70**, 1189 (1997).

[16] S. Hughes and D. S. Citrin, Phys. Rev. B **59**, 5288 (1999).

[17] D. S. Citrin, Phys. Rev. B **47**, 3832 (1993).

[18] For how the tunneling Hamiltonian leads to the Green functions via a scattering theory and the relation with transfer matrices, see Refs. 4,5,8 and N. S. Wingreen *et al.*, Phys. Rev. B **40**, 11 834 (1989); A. D. Stone and P. A. Lee, Phys. Rev. Lett. **54**, 1196 (1985); Y. G. Rubo, Zh. Éksp. Teor. Fiz. **104**, 3536 (1993) [JETP **77**, 685 (1993)]; B. Y. Gelfand *et al.*, Phys. Rev. Lett. **62**, 1683 (1989).

[19] F. L. Lederman and J. D. Dow, Phys. Rev. B **13**, 1633 (1976); D. A. B. Miller *et al.*, *ibid.* **32**, 1043 (1985); D. A. B. Miller *et al.*, IEEE J. Quantum Electron. **QE-22**, 1816 (1986).

[20] G. N. Watson, *A Treatise on the Theory of Bessel Functions* (Cambridge University Press, Cambridge, England, 1944).

[21] M. S. C. Luo *et al.*, Phys. Rev. B **48**, 11 043 (1993); A. P. Heberle *et al.*, Phys. Rev. Lett. **75**, 2598 (1995).

[22] D. S. Citrin and T. B. Norris, IEEE J. Sel. Top. Quantum Electron. **2**, 401 (1996).

Session IC15 - THz-Spectroscopy and Ultra Fast Optics.
FOCUS session, Tuesday morning, March 23
Room 265W, GWCC

[IC15.07] Ultrafast coherent effects in semiconductor quantum wells using modulated optical fields

W Harshawardhan, S Hughes, D. S. Citrin (Washington State University)

We theoretically study ultrafast coherent dynamics in quantum wells, excited by broadband optical pulses that may be frequency or amplitude modulated. Both excitonic and free-carrier effects are included self-consistently by numerically solving the semiconductor Bloch equations. Signatures of coherent population trapping and adiabatic population transfer between the 1s exciton and the continuum are demonstrated. Other novel dynamical effects will be discussed.

Part I of program listing

PHYSICAL REVIEW B VOLUME 60, NUMBER 23 15 DECEMBER 1999-I

Excitonic-state trapping and quasiadiabatic population transfer in a two-band semiconductor

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By solution of the semiconductor Bloch equations in the presence of a frequency-modulated optical pulse, clear exciton trapping and quasiadiabatic population transfer in a two-band semiconductor is predicted. [S0163-1829(99)00647-5]

Driven by the pursuit to understand the differences and similarities between the semiconductor and atomic systems resonantly excited by laser pulses, the coherent nonlinear dynamics of direct-gap semiconductors has been vigorously investigated over the years. Moreover, with the ongoing advancement of short-pulse laser techniques and excellent semiconductors samples, classes of coherent dynamic phenomena have been discovered recently including, for example, coherent exciton control[1,2] and self-induced transmission.[3] Running in parallel have been discoveries in atomic optics such as population trapping using frequency modulated fields. A realization of such state trapping and its periodic redistribution in a two-level atom (TLA) was predicted in Ref. 4. In the past year, in combination with other well-known concepts such as adiabatic rapid passage, multiphoton resonances, and Landau-Zener transitions, trapping was experimentally demonstrated in a TLA [5]. For two-band semiconductors, however, trapping is scarcely expected due to Coulomb many-body complications and the valence/conduction band-continua of free-carriers. Surprisingly in this work we predict clear excitonic-state trapping (EST) and quasiadiabatic population transfer (QAPT) using multicolored excitation schemes.

For high-intensity ultrashort optical pulses, it is well established that Rabi oscillations of the population between two states can be seen in the temporal evolution of a TLA.[7] However, in the case of a sinusoidal frequency-modulated excitation, square-wave oscillations of the population—periodic state trapping—as well as more complicated and phase-dependent structures may appear. In semiconductors, the *two-level* model (as a first approximation to a two-band description) is considered inappropriate because Coulomb many-body interactions result in a renormalized Rabi energy and bandedge, and excitation-induced dephasing (EID). Nevertheless, self-induced transmission and multiple Rabi flopping on a semiconductor free exciton resonance have been recently reported in bulk[3] and quantum wells (QW's);[6] these measurements were successfully explained within the framework of the semiconductor Bloch equations. Naturally the question arises: Can one achieve population trapping dynamics in semiconductors using frequency-modulated or suitably-chirped broadband optical pulses? By employing, for example, a frequency-dependent 150-fs full width at half maximum (FWHM) irradiance sech pulse excited at the $1s$ exciton peak, we predict, *yes*, one can achieve EST and QAPT in a semiconductor, too. We present trapping dynamics in semiconductors akin to trapping in atomic systems.[4] We discuss in detail conditions to achieve EST and QAPT even when a broadband pulse excites both Coulomb-bound electron-hole (e-h) pairs (excitons) and free e-h pairs. The other well known trapping phenomenon in atomic systems is the coherent population trapping effect, which has been widely used *inter alia* in adiabatic population transfer, electromagnetically induced transparency, and velocity selective cooling. As pointed out recently by using adiabatic population transfer between heavy- and light-hole bands (to mimic a three-level atom),[8] the exact analog to the coherent-population-trapped state (dark state) used in atomic physics is not possible in semiconductors.

For our theoretical approach we assume the validity of the rotating-wave approximation and thus neglect the possibility of carrier-wave Rabi flopping.[9] We further assume a two-band QW where each e-h state within a certain bandstructure with a wave number $\mathbf{k}$, contributes to the total optical polarization $P=2A^{-1}\Sigma_{\mathbf{k}}d_{\mathrm{cv}}P_{\mathbf{k}}$, with d_{cv} the transition dipole matrix element. The SBE for the polarization functions are ($\hbar=1$) (Ref. 10) $\dot{P}_{\mathbf{k}}=-i\Delta_{\mathbf{k}}P_{\mathbf{k}}-i\Omega_{\mathbf{k}}(f_{\mathbf{k}}^{e}+f_{\mathbf{k}}^{h}-1)+\dot{P}_{\mathbf{k}}|_{\mathrm{corr}}$, where $\Delta_{\mathbf{k}}$ is the renormalized energy dispersion for a parabolic two-band semiconductor and $\Omega_{\mathbf{k}}$ is the generalized Rabi frequency, while $f_{\mathbf{k}}^{e/h}$ are the carrier distribution functions of the electrons and holes. The Coulomb interaction is treated here in a quasi-static approximation (single plasmon-pole).[10] Similarly, the carrier density can be calculated from the electron (or hole) population $N=2A^{-1}\Sigma_{\mathbf{k}}f_{\mathbf{k}}$, where $f_{\mathbf{k}}$ is calculated from $\dot{f}_{\mathbf{k}}^{e/h}=iP_{\mathbf{k}}^{*}\Omega_{\mathbf{k}}-iP_{\mathbf{k}}\Omega_{\mathbf{k}}^{*}+\dot{f}_{\mathbf{k}}^{e/h}|_{\mathrm{corr}}$. In addition to the terms that result from the time-dependent Hartree-Fock approximation, the carrier-carrier (CC) and carrier-phonon (CP) collisions drive the nonequilibrium distribution functions towards quasiequilibrium Fermi functions and yield optical dephasing. The CC collisions (Coulomb correlation terms) are calculated from the e-h Boltzmann equations including also a correlation field, nondiagonal dephasing, and polarization scattering.[11–13] The influence of CP interactions occurs over much longer time scales than those studied in this work and can be safely neglected.

We assume input optical pulses of the form $\mathbf{E}(\mathbf{r},t)=E(t)e^{-i[\omega_0 t+k_0 z+\phi(t)]}+\mathrm{c.c.}$, polarized in the plane of the QW, with all material parameters corresponding close to $In_xGa_{1-x}As/GaAs$ QW's (Ref. 14) and $\omega_0=\omega_x$ ($1s$ exciton resonance). $In_xGa_{1-x}As/GaAs$ QW's are advantageous since with compressive strain one can increase the splitting of the heavy- and light-hole exciton and thus neglect the light-hole states, therefore validating the two-band model. In analogy with the atomic systems we modulate the input field by the phase factor $\phi(t)=M\sin(\Omega t)$, with M and Ω the index and frequency of modulation. For a TLA the resonant part of the

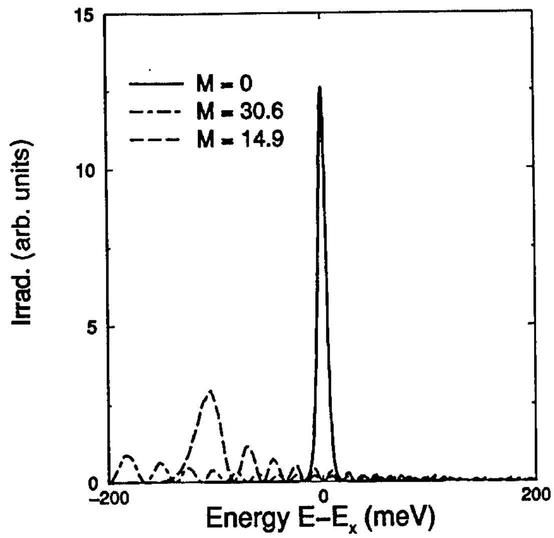

FIG. 1. Irradiance spectra for a frequency-modulated and unmodulated 150 fs optical pulse. Energy is in units of the exciton binding frequency.

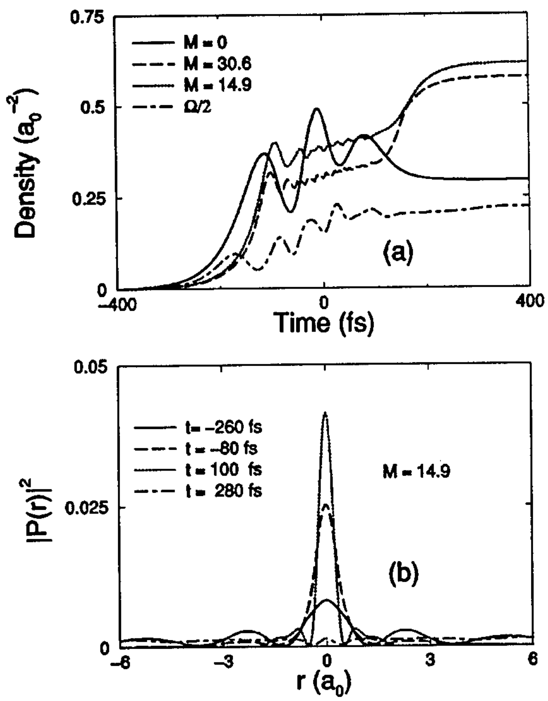

FIG. 2. (a) Pulse-induced carrier density (in units of inverse Bohr radius squared) as a function of time. The chain line corresponds to $M=14.9$ but with a frequency modulation of $\Omega/2$. (b) The polarization density corresponding to the dashed case in Fig. 2(a) at several snapshots.

interaction Hamiltonian vanishes if M is chosen such that the bessel function $J_0(M)=0$ and one is left with the nonresonant rapidly oscillating terms.[4] The solid-state community may recall that a somewhat similar criterion has been derived to realize dynamic localization in infinite lattices driven by a harmonic time-dependent electric field.[15,16]

In the TLA it is further known that the two bare levels cross at the times $t_n = n\pi/2\Omega$ ($n=$ integer). The crossing of energy levels occurs quite naturally in the instantaneous frame of the modulated field. The population distribution at the crossings can be understood semiquantitatively by integrating numerically the time-dependent Schrödinger equation or by employing Landau-Zener theory.[4] For semiconductor the situation is much more complex, whereby the optical field provides coupling between two bands, and a modulated near-resonant field leads to coupling of higher spectral components of the field to larger **k** states in the band, and the field components below the resonance provide a nonresonant coupling to the excitonic transition. Moreover, the Coulomb interaction yields exciton and plasma induced many-body effects. These conditions prohibit us from having a simple set of criteria for trapping. Thus we propose here a novel spectral distribution for the field, which has its central frequency detuned far below the excitonic resonance, and, by strategic sweeping of the instantaneous carrier frequency, dynamical population trapping can be realized. As in the TLA case we do not necessarily have to choose M to be strictly a zero of the Bessel function.

For the modulation frequency we choose $\Omega=8$ meV unless stated otherwise. As mentioned above, for the *unmodulated* field, $E(t)$, we employ a 150-fs pulse excited at the $1s$ exciton peak. In Fig. 1, we depict the spectral irradiance of the unmodulated and frequency-modulated fields. With modulation, a much larger bandwidth can be obtained with a series of larger peaks extending far below the exciton resonance; the injected pulse profile is sufficiently broad to couple many different excitation modes. One recognizes that the dominant spectral components are well below the band edge. By increasing the index of modulation (M), the spectral width of the pulse increases. In contrast, without any modulation the spectrum of the excitation pulse is significantly narrower and peaked at the central frequency as shown in Fig. 4(b) (solid line; see later). Although such a modulation is difficult to obtain experimentally, one only requires about one approximate oscillation to obtain the appropriate energy level crossings. Indeed one can also achieve trapping phenomena using suitably chosen linear chirps (discussed below). Further, recent developments of free-electron lasers as well as terahertz solid-state emitters may be utilized in materials to achieve the desired frequency modulation in a suitably chosen nonlinear material.[17]

For an optical pulse that has roughly an area (integral of its Rabi frequency over time) of 6π (~ 1 GW/cm^2), Fig. 2(a) shows the excitation-induced density for the QW optically excited at the $1s$ exciton resonance using (i) an unmodulated input pulse [the solid curve depicts the carrier-density showing the familiar Rabi flopping[7,18] of the density (complete inversion is not possible due to Coulomb processes)]; and (ii) a frequency-modulation of the input pulse for two different values of the index of modulation M: (the dashed and dotted curves are markedly different from the unmodulated case and display strong EST and ultrafast QAPT). The values of M were chosen in accord with the original work in atoms, while the pulse area was chosen to be large enough to drive the population into trapping–it is also close to values employed experimentally.[6] During the time interval from about -160 to 160 fs, the population remains trapped in the exciton level, eventually jumping out in the wake of the exciton-continuum quasi-adiabatic crossing. For the case of $\Omega/2$ and $M=14.9$ (chain curve), the trapping efficiency is much less as expected and QAPT does not occur, since the crossing occurs just once and the on-resonance field components are not sufficient to create the trappinglike feature; however, even in a regime where the pulse irradiance is negligible there is some evidence for small density changes at around 360 fs. In all cases, there are signatures of phase interference due to coherent carrier evolution along different interfering pathways and EID. We would like to point out that to obtain such good Rabi flopping one must include EID at a microscopic level that includes both nondiagonal dephasing and polarization scattering;[11] these contributions from CC scattering reduce the interband optical linewidths of the higher $\mathbf{k}$ states and thus limit the higher-energy continuum occupations in comparison to the pure dephasing (essentially a relaxation-time approximation) treatment.[12] Physically, this is important since the leading edge of the pulse, which is detuned below the exciton resonance, prepares the system for the excitation of real population. A state and energy-independent dephasing time would result in erroneous large dephasing of the initially virtual excitation and therefore suppress the QAPT and EST. Polarization scattering is also known to result in a transfer of oscillator strength from the continuum to the exciton.[13]

To highlight the difference between the Rabi flopping and trapping we depict in Fig. 2(b) the polarization density (wave-packet), $|P(\mathbf{r},t)|^2$ with $\mathbf{r}=\mathbf{r}_e-\mathbf{r}_h$ at various temporal snapshots.[19] During the period -260 to 100 fs the population becomes strongly trapped in the exciton state indicated in the figure by a high probability of finding the electron and hole at the same relative position. However at the later time of 280 fs, after the density increases rapidly (we mention that this increase is almost steplike in the absence of dephasing), the excitonic probability *decreases* substantially and the wave packet speads out significantly, demonstrating that the population is no longer trapped. This drammatic spreading of the wave packet arises due to the modulation and the resulting crossing of carriers into the continuum. The above picture sheds much more light on the trapping scenario than, for example, $f_{\mathbf{k}}^{e,h}$ or $\mathbf{k}$, which do not clearly distinguish between Coulomb-bound excitons and free carriers in the continuum. The spatial polarization dynamics at $\mathbf{r}=0$ may in fact be probed experimentally using conventional four-wave mixing techniques. We observe here a method to control, coherently, the excitonic wave packet by tailoring the conditions of the energy level crossing. To probe the entire spatial dynamics one would need to, e.g., couple a THz field with the optical field; this would involve a significantly more complex analysis. We also mention that to direcly detect density changes, experimentally pulse-propagation studies are very difficult and one should, for example ,use the technique reported recently in Ref. 6. This allows for the detection of density changes via simple differential transmission changes of a probe pulse.

The aforesaid phenomena beg a more physically intuitive explanation: Adding a frequency-modulation to the excitation field is equivalent to modulating the energy separation between the ground and exciton/continuum states. In Fig. 3 we show, schematically, the various crossings of energy levels that lead to quasiadiabatic transfer of population at these crossings. The crossings "$B-D$" transform into anticrossings (solid lines) due to the coupling with the effective Rabi field. In contrast the crossing at "A" is unaffected due

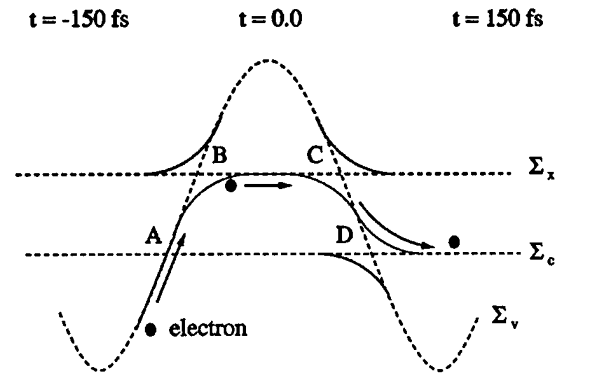

FIG. 3. Schematic of the energy level crossing that results from modulation of the field, where "$A-D$" represent the various crossings. The arrows depict the evolution of the population (large black dots).

to the weak coupling of the $v\leftrightarrow c$ transition on the leading edge of the pulse. The modulation of the dressed energy levels results from transforming into a frame corresponding to the instantaneous field frequency, and are given as $\Sigma_c=-[\Delta_{\mathbf{k}}^2-\Delta_x^1]$, $\Sigma_x=0$ and $\Sigma_v=-[\Delta_x^1+M\Omega\cos(\Omega t)]$, for the electrons in the conduction band (c), excitonic state (x), and electrons in the valence band (v), respectively. One can of course depict the diagram in reverse for the holes. Here, Δ_x^1 ($\Delta_{\mathbf{k}}^2$) denote the detuning of the central laser field frequency from the excitonic (continuum) energy levels. The coupling of the field to the $v\leftrightarrow x$ and $v\leftrightarrow c$ transitions transforms these crossings into *avoided crossings*. Initially the system will be off resonance and far from any crossing; as the modulation changes, the population is swept though resonance and it evolves quasi-adiabatically into the excitonic state where it displays the trapping feature. The system further encounters two closely spaced crossings, resulting in an enhanced step-like transfer of population into the continuum. The arrows in Fig. 3 indicate one possible temporal path (high probability) along which the electrons may evolve ($\Delta_{\mathbf{k}}^2$ is fixed for simplicity). The detuning term $\Delta_{\mathbf{k}}^2$-Δ_x^1 is the difference between the continuum energy levels and frequency of the the $1s$ exciton peak. One should keep in mind however that the above model is grossly simplified and many-body effects, included in our numerical results, will complicate things substantially; however our quantitative theoretical study is in fairly good agreement with the above level crossing model.

The essential difference in the trapping criterion between the atomic case and semiconductor case lies in the frequency content of the exciting field. In the atomic case the trapping-like phenomenon with nearly complete inversion results from *correlated sideband excitation* of the atom.[20] The frequency components of the modulated field excite the atom symmetrically about the atomic resonance, resulting in trapping. In the semiconductor, a symmetric frequency content of the excitation does not lead to the desired trapping due to its nonsymmetric coupling to the band structure of the semiconductor; effectively the high-frequency components of the modulated pulse selectively excite the continuum of states resulting in large carrier generation, thus washing out the trapping. We circumvent this problem by employing a modulated pulse with its predominant frequency content away from the band edge.

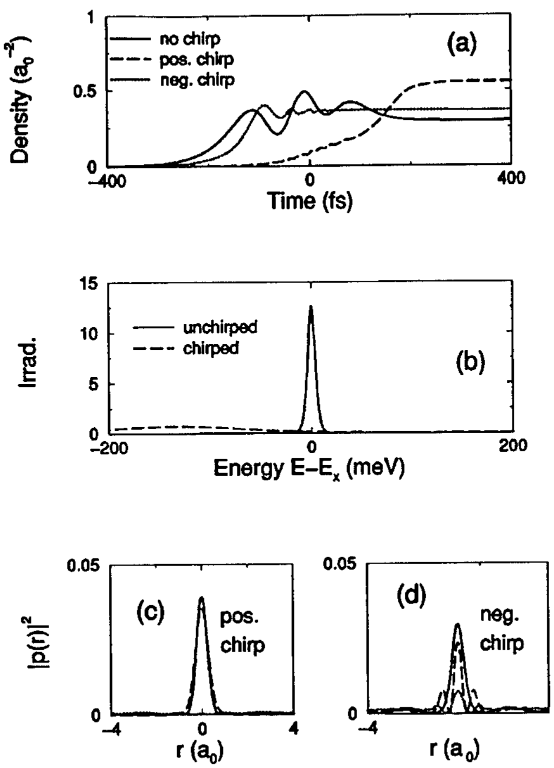

FIG. 4. (a) Pulse-induced carrier density as a function of. (b) Corresponding input pulse spectra for (a). (c) The polarization density for the positive chirp corresponding to the dashed case in (a) at several snapshots (see text). (d) As in (c) but for negative chirp.

Next, we employ linear-chirped input pulses to discern if some of the same qualitative trapping features (EST and QAPT) can also be obtained using suitably chosen chirp parameters. To best immitate the frequency-modulated spectrum, we choose a frequency chirp of the form: $\omega \rightarrow \omega_x - 120$ meV$+a_c t$ with a_c being ± 0.5 meV/fs. $E_0(t)$ is the same as before. The chirped-pulse spectrum along with the unchirped spectrum is displayed in Fig. 4(b). The pulse-induced density with positive a_c is shown by the dashed line which, although does not show signs of trapping, does show fast oscillations in the density as well as a rapid increase in the density near the crossing time. For the negative a_c, strong excitonic trapping is again achieved and seems to maintain its trapped state since no further crossing takes place. This is again clear from the simplified zero-order energy-level picture with appropriate crossing of the $1s$ exciton state, ground crystal state, and continuum states. The slope and sign of the chirp (a_c) gives us a handle to selectively suppress or enhance the excitonic transition. The trapping and population transfer are highlighted in Figs. 4(c)–4(d) which show the wave-packets at the times -80 fs (solid curve), 100 fs (dashed curve), and 280 fs (dotted curve) for both the negative and positive chirp pulse excitation. Population transfer is not expected for the negative-chirp case as the zero-order energy level never cross.

In conclusion, we predict the possibility of exciton trapping and quasiadiabatic population transfer in a two-band semiconductor using frequency-modulated optical pulses. It is clear that the trapping feature arises out of an interplay of the dominant excitonic resonance and the excitation by a sufficiently *strong, broadband off-resonant pulse*, with a weak yet broad spectral content to excite the continuum of states. The population redistribution results from the crossing of energy levels. Besides being an intriguing theoretical study, our results are timely with recent advances in frequency-modulated spectroscopy techniques and the observation of multiple Rabi flopping on free exciton transitions.

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[1] A. P. Heberly *et al.*, Phys. Rev. Lett. **77**, 2598 (1995).

[2] D. S. Citrin, Phys. Rev. Lett. **77**, 4596 (1995).

[3] H. Giessen *et al.*, Phys. Rev. Lett. **81**, 4260 (1998).

[4] G. S. Agarwal and W. Harshawardhan, Phys. Rev. A **50**, 4465 (1994).

[5] M. W. Noel *et al.*, Phys. Rev. A **58**, 2265 (1998).

[6] A. Schülzgen *et al.*, Phys. Rev. Lett. **82**, 2346 (1999).

[7] See, for example, L. Allen and J. H. Eberly, *Optical Resonance and Two-Level Atoms* (Wiley, New York, 1995).

[8] R. Binder and M. Lindberg, Phys. Rev. Lett. **81**, 1477 (1998).

[9] S. Hughes, Phys. Rev. Lett. **81**, 3363 (1998).

[10] M. Lindberg and S.W. Koch, Phys. Rev. B **38**, 3342 (1988).

[11] A. Knorr *et al.*, Chem. Phys. **210**, 27 (1996).

[12] S. Hughes *et al.*, J. Opt. Soc. Am. B **49**, 754 (1997).

[13] T. Rappen *et al.*, Phys. Rev. B **49**, 10 774 (1994).

[14] Reduced mass, $\mu = 0.035\ m_e$; $1s$ excitonic binding energy, $E_{1s} \approx 12$ meV; and Bohr radius (2D), $a_0 = 75$ Å.

[15] D. H. Dunlap and V. M. Kenkre, Phys. Rev. B **34**, 3625 (1986).

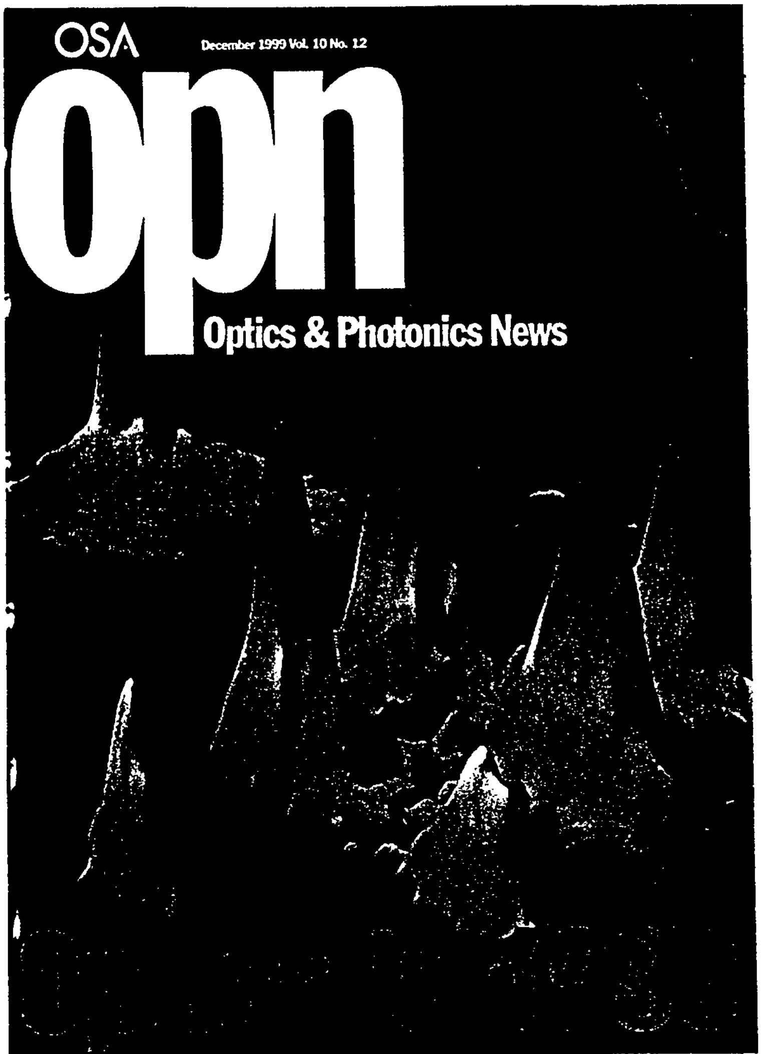
[16] M. Holthaus, Phys. Rev. Lett. **69**, 351 (1992).

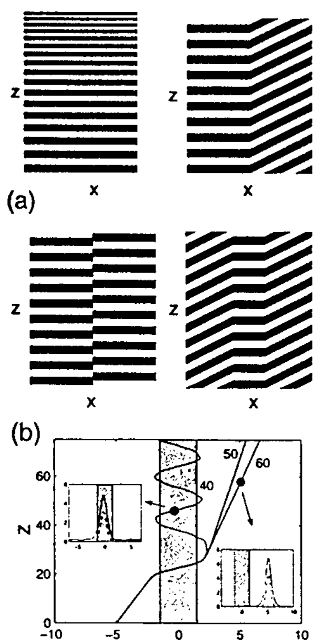
[17] G. von Plessen (private communication).

[18] I.I. Rabi, Phys. Rev. **51**, 652 (1937).

[19] W. Harshawardhan and G. S. Agarwal, Phys. Rev. A **55**, 2165 (1997).

[20] P.K. Lam and C.M. Savage, Phys. Rev. A **50**, 3500 (1994).





Torner Figure 1. Sketch of possible QPM lay-outs with chirps, dislocations, tilts and wells, and illustration of power-dependent trapping of solitons in a QPM well. Labels in **(b)** stand for different soliton input energies.

expanding further the commercial applications of QPM devices.[2] However, QPM engineering finds novel important applications far beyond pure frequency-conversion devices.

We have recently exposed the great promise of engineered QPM patterns in soliton systems. Solitons in QPM structures form by the mutual parametric trapping of the multi-frequency light waves that propagate in the crystal.[3] The "glue" that keeps the multi-frequency signals locked and trapped in robust soliton light states is their energy- and phase-exchanges, mediated by cascading of quadratic nonlinearities. Therefore, all the stationary and dynamic features of the soliton signals depend crucially on the local phase-mismatch experienced by the multi-frequency waves that form the solitons.

The breakthrough introduced by QPM engineering is the possibility to tailor the longitudinal and transverse variations of the phase-mismatch between the waves, hence to tailor the properties of the solitons. This is a unique possibility that links the implementation of potential soliton devices with the identification and design of innovative topologies of QPM patterns. In particular, we have shown the soliton based signal compression and general shaping, which includes not only tailoring the soliton width and profile but also the fraction of energy carried by each of the waves forming the multi-color signals, in QPM structures with engineered longitudinal chirps;[4] and the spatial soliton switching in QPM settings with engineered transverse variations of the phase mismatch.[5] Transverse patterns illustrative of such possibilities include chirps, tilts, dislocations and wells, as those shown in Figure 1a. In the particular example shown in Figure 1b, upon strong illumination the soliton crosses the QPM well. However, with a weaker illumination the soliton is trapped in the well. Soliton self-bouncing and power-dependent soliton scanning are other potential feasible possibilities shown. More advanced and practical QPM lay-outs, including modulated, multi-periodic, and quasi-periodic gratings, and engineered QPM multi-grating arrays are now under investigation. Engineered competing quadratic and cubic effective nonlinearities,[6] bandwidth-tailored soliton excitation, quasi-discrete soliton signals, and novel kinds of stable solitons constitute important examples of the fascinating potential outcome.

We believe that these findings constitute the tip of the iceberg of the potential of QPM engineering for quadratic soliton control devices in single-, multi-pass, and cavity geometries. We anticipate that creative QPM engineering might offer crucial building blocks of all future photonic devices employing quadratic soliton signals.

References

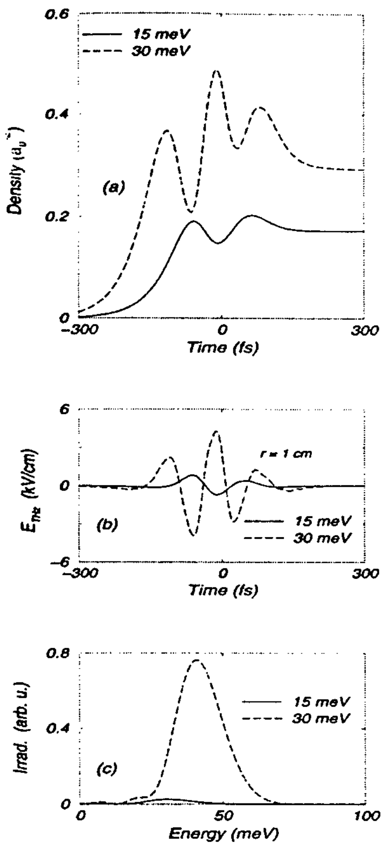
1. M.M. Fejer, " Nonlinear optical frequency conversion," *Phys. Today*, **46** (5), 25–32 (1994).
2. G. Sucha, "Doped optical fibers promise compact femtosecond sources," *Laser Focus World* **34** (4), 121–6 (1998); S. Bains, "PPLN inspires new applications," *Laser Focus World*, **34** (5), 16–19 (1998).
3. L. Torner, *et al.*, "Soliton evolution in quasi-phase-matched second-harmonic generation," *J. Opt. Soc. Am. B*, **14**, 3127–33 (1997); C.B. Clausen *et al.*, "Spatial solitons and induced Kerr effects in quasi-phase-matched quadratic media," *Phys. Rev. Lett.*, **78**, 4749–52 (1997).
4. L. Torner, *et al.*, "Adiabatic shaping of quadratic solitons," *Opt. Lett.*, **23**, 903–5 (1998).
5. C.B. Clausen, *et al.*, "Spatial switching of quadratic solitons in engineered quasi-phase-matched structures," *Opt. Lett.*, **24**, 7–9 (1999).
6. O. Bang, *et al.*, "Engineering competing nonlinearities," *Opt. Lett.*, **24**, 1413–15 (1999).

Terahertz Sources

Extremely-Tunable Terahertz Emission: Coherent Population Flopping in a DC-Biased Quantum Well

S. Hughes and D.S. Citrin

Terahertz (THz) light-emitters are a welcome addition to the laser engineer's palette since they offer a marriage of optics and high-speed electronics. In the field of semiconductor nanostructures, the exploration of THz generation schemes and THz-induced carrier dynamics provides a useful means to study charge-carrier transport. The inter-



Hughes Figure 1. (a) Temporal variation of carrier pulse-induced carrier density in units of inverse Bohr radius squared; excitation is at the 1*s* exciton resonance. The peak amplitude of the pulse is 15 meV (solid curve) and 30 meV (dashed curve). **(b)** Emitted THzfield versus time for the 15 meV and 30 meV pulses at $r = 1$ cm from the excitation spot. **(c)** Irradiance of the emitted THz spectra corresponding to **(b)**.

play of optical and THz transients, moreover, provides a unique window for simultaneously studying inter- and intra-band carrier dynamics in semiconductor heterostructures. Applications with the semiconductor laser[1] include gain modulation and all-optical ultrafast switching rates of around 500 gigabit/sec—a useful prerequisite for the next generation of broad-bandwidth optical communications networks and the inevitable *gigabit Internet.* More fundamental applications encompass coherent control of electronic wavepackets,[2] nanotechnology, and quantum computing. In short: the cry for reliable THz emitters can be heard throughout the globe. However, apart from expensive free-electron lasers, a common limitation to the generation of THz transients in the laboratory is the lack of tunable sources.

The successful advancement in short-pulse laser techniques and excellent semiconductors samples led to a breakthrough in the last year by the clear observation of self-induced transmission (SIT) in bulk semiconductors[3] and multiple Rabi floppping in quantum wells (QWs).[4] Rabi flopping is a well known coherent effect in atomic systems but has long been considered to be the *Holy Grail* to the *ultrafast* semiconductor community—who are more than ever driven by the pursuit to understand the differences and similarities between the semiconductor and atomic systems resonantly excited by laser pulses. Very recently we have devised a new method[5] that exploits Rabi-flopping-induced density oscillations in a weakly-DC-biased QW to realize very intense (~kV's/cm), tunable terahertz transients. In our scheme, extreme— *tailored*—tunability is achieved by simply varying the area of the input optical pulse. We employ a 150 fs FWHM irradiance pulse. For a weakly-biased QW we assume $d = 0.8\ e\text{Å}$(corresponding bias field ~ 10–100 V/cm), and calculate the emitted field at $r = 1$ cm. We assume the validity of the slowly varying envelope approximation (which is not always valid, see Reference 6), and compute the semiconductor Bloch equations including all the relevant scattering mechanisms.

Figure 1a shows the excitation-induced density for the QW optically excited at the 1*s* exciton resonance using the input Rabi energies of 15 meV and 30 meV (pulse irradiance: ~ GW/cm^2), respectively. The curves display the conduction band-valence band Rabi flopping[3,4] (complete inversion is not possible due to Coulomb processes). Material parameters correspond to $In_xGa_{1-x}As$/GaAs QW's and agree well with the experimental observations. In Figure 1b we show the emitted THz field versus time corresponding to the density oscillations in Figure 1a. Each transient is approximately 500 fs in duration (due to dephasing) and undergoes a series of polarity changes. In the EM spectra (Figure 1c), clear peaks in the THz regime are seen. For larger areas, more Rabi flops occur[3,5] that subsequenly lead to THz emission with larger frequencies. We note that current techniques employing sub-ps optical pulses for THz-generation are usually limited to about the 1–3 THz regime at best. Our selected examples here easily go to the 10s-of-THz regime with the added bonus of much larger fields. Futhermore they do not suffer from any bandwidth limitations (c.f. optical rectification techniques). These studies are motivated in part because of recent dramatic advances in high-speed, hybrid lightwave-THz photonic applications.

Acknowledgments

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References

1. S. Hughes and D.S. Citrin, "Ultrafast heating and switching of a semiconductor optical amplifier using half-cycle terahertz pulses," *Phys. Rev. B*, **58**, R15969 (1998); C.Z. Ning, *et al.*, "Ultrafast modulation of semiconductor lasers through a terahertz field," *Appl. Phys. Lett.*, **75**, 442 (1999).
2. S. Hughes and D.S. Citrin, "Creation of highly anisotropic wavepackets in quantum wells: dynamical Franz-Keldysh effect in the optical and terahertz regimes," *Phys. Rev. B.*, **59**, R5288 (1999).
3. H. Giessen *et al.*, "Self-induced transmission on a free exciton resonance in a semiconductor," *Phys. Rev. Lett.*, **81**, 4260 (1998).
4. A. Schulzgen *et al.*, "Direct observation of excitonic Rabi oscillations in semiconductors," *Phys. Rev. Lett.*, **11**, 2346 (1999).
5. S. Hughes and D. S. Citrin, "Tunable THz emission via multiple Rabi flopping in a DC-biased quantum well: a new strategy," *Opt. Lett.*, **24**, 1242 (1999)
6. S. Hughes, "Breakdown of the area theorem: carrier-wave Rabi flopping of femtosecond optical pulses," *Phys. Rev. Lett.*, **81**, 3363 (1998).

Ultrafast Lasers

Real-Time Optical Pulse Characterization Using SPIDER

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The recent advances in ultrafast lasers have mandated a change in the techniques used to measure them. It is no longer sufficient to measure merely the spectrum or the auto-correlation of an ultrafast laser pulse; a more complete picture is required, and this involves determining the phase of the time-varying optical field. The ability to rapidly change the shape of the pulse is a key technolo-